Photodisintegration of the deuteron in the few GeV region using asymptotic amplitudes

A.E.L.Dieperink[a], S.I.Nagorny[b]

[a] *Kernfysisch Versneller Instituut, NL-9747 AA Groningen*

[b] *Nikhef, P.O.Box 41882, NL-1009 DB Amsterdam, The Netherlands*

Abstract

Exclusive photodisintegration of the deuteron in the 1-4GeV range is described in terms of a simple covariant and gauge invariant approach using an effective counting rule as the hard part of the d-np vertex. At a scattering angle of $\theta_{cm} = 90°$ a prescaling behavior of the differential cross section $\propto s^{-(n-2)}$ with $n \approx 12$ is obtained; going away from 90° the value of n decreases slowly, in qualitative agreement with the recent data.

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I. INTRODUCTION

The exclusive process of photodisintegration of the deuteron addresses an interesting interplay of nuclear and particle physics. At low energies (say, below $E_\gamma = 0.5$ GeV) conventional nuclear models based upon meson exchange which fit the NN phase shifts give a satisfactory description of both the energy dependence and the angular distribution of the experimental cross section [1–3]. However, at higher energies nuclear potential models fail to explain the data [4,5]. This is not unexpected since at $E_\gamma > 1$ GeV small distances of the order of 0.2 fm play a role.

Alternatively it has been attempted to describe the cross section in terms of quark-gluon

degrees of freedom [6–8]. A possible signature for the emergence of quark-gluon degrees of freedom would be the observation of the onset of scaling of the cross section [6]. Several examples where this happens have been found. For example, the $p(\gamma,\pi)n$ reaction above 3 GeV appears to be consistent with counting rules at all angles [9].

In case of the exclusive process $d(\gamma,p)n$ the dimensional counting analysis of [6] leads to a differential cross section of the form

$$d\sigma/dt = s^{-(n-2)} f(\theta), \tag{1}$$

where $s(\theta)$ is the cm energy (angle), and n denotes the number of elementary fields in initial and final state (i.e. $n = 13$ in case of the deuteron).

Previous data from SLAC [4] up to E_γ =3 GeV at $\theta = 90^\circ$ indicated a scaling behavior consistent with $n = 13.1\pm0.3$. More recent data from Jefferson Lab up to 4GeV confirmed [5] the scaling behavior with $n \approx 13$ at $\theta = 90^\circ$ and 69°; whereas at smaller angles, $\theta = 36^\circ, 52^\circ$, the best fit yielded $n \approx 11.5$ and 11.6, respectively. This constitutes a deviation from the scaling behavior predicted by simple counting rules.

A more refined approach is the reduced nuclear amplitude (RNA) approach [7] which is also based upon parton exchange between the two nucleons, but takes into account some finite mass and higher twist effects. However, if normalized at $E_\gamma = 1.0$GeV, the prediction falls below the data at $\theta = 90^\circ$ for $E_\gamma > 3$GeV. On the other hand the quark-gluon string (QGS) model proposed in [8] and based upon Regge phenomenology appears to describe the data [5] only at small t values, corresponding to small angles.

This indicates that in practice the situation is more complex.

The aim of this paper is to study the question of the possible origin of the apparent scaling and scaling violation in more detail; in particular we address whether the occurrence of scaling is an exclusive pQCD phenomenom or whether it can arise from different mechanisms. The approach in the present paper, in which the basic degrees of freedom are taken to be hadronic, is an extension of the one in ref. [10] which predicted the cross section at $\theta = 90^\circ$ for $E_\gamma > 1$ GeV in fair agreement with experiment, but failed to describe the angle

dependence. Using a simple covariant parametrization of the deuteron vertex in terms of a hard component (suggested by an effective counting rule) and imposing gauge invariance the cross section for large but not infinite s shows a "preasymptotic" scaling behavior at $\theta = 90°$. This resembles the counting rule prediction, however, with an angle dependent value of n (which decreases away from 90°).

II. FORMALISM

The general covariant half off-shell d-pn vertex ($p_2^2 = m^2$) can be expressed as [12]

$$A^\mu = \Gamma_1 \gamma^\mu + \frac{(p_1 - p_2)^\mu}{2m}\Gamma_2 + \frac{\not{p}_1 - m}{2m}[\Gamma_3 \gamma^\mu + \frac{(p_1 - p_2)^\mu}{2m}\Gamma_4], \tag{2}$$

where the Γ_i are scalar functions: invariant form factors. For large virtualities the contribution of the first term in eq.(2), proportional to Γ_1, dominates over the last three. Assuming that at large s only tree-type diagrams survive, the $\gamma + d \to n + p$ (Born) amplitude can be written as the sum of the pole diagrams in the s, t and u channels ($\Gamma_1 \equiv G$):

$$T^\mu_{pole} = \xi_\nu(d)\bar{u}(p')[F^p_\mu \frac{\not{p} + m}{p^2 - m^2}\gamma^\nu G(-k_t^2) + F^n_\mu \gamma^\nu \frac{\not{n} - m}{n^2 - m^2} G(-k_u^2)$$

$$+ G(-k_s^2)\gamma^\alpha \frac{-g_{\beta\alpha} + d'_\beta d'_\alpha / d'^2}{s - m_d^2} F^d_{\mu\nu\beta}] C \bar{u}^T(n'). \tag{3}$$

Here s-,t-, u- variables for the d-pn vertex are: $k_t = (k_s - q)/2,\ k_u = (k_s + q)/2,\ k_s = (p' - n')/2,\ p' = p + q,\ d' = d + q$; $u(p)$ is the nucleon spinor, $\xi_\nu(d)$ the polarization vector of the deuteron ($\xi^\nu(d) d_\nu = 0$), C is the charge conjugation operator, F^i_μ ($i = n, p$) denotes the electromagnetic coupling to the nucleon, $F^i_\mu = (e_i + \not{q}\frac{K_i}{2m})\gamma^\mu$, and $F^d_{\mu\nu\beta}$ the corresponding one for the deuteron [10] ($f_d = 2\mu_d - 1 + Q_d$):

$$-F^d_{\mu\alpha\beta} = 2d_\mu(g_{\alpha\beta} - q_\alpha q_\beta \frac{f_d}{2m_d^2}) + 2\mu_d(g_{\mu\alpha} q_\beta - g_{\mu\beta} q_\alpha) + (s - m_d^2)(g_{\mu\alpha} q_\beta + g_{\mu\beta} q_\alpha)\frac{f_d}{4m_d^2}.$$

Both γNN and γdd vertices satisfy the identities: $q^\mu F^i_\mu = e_i \not{q}$, $q^\mu F^d_{\mu\alpha\beta} = -(s - m_d^2) g_{\alpha\beta}$. Note, s-channel accounts for the pole part of the T-matrix of the final state n-p interaction. In the presence of a momentum dependent vertex function the pole diagrams themselves are

not gauge invariant. Indeed, using the Ward-Takahashi identity for the 3-point electromagnetic vertices, eq.(3), and the Dirac equation one finds that the contraction of q_μ with the sum of the s-, t-, u-pole amplitudes does not vanish (i.e. the corresponding Born current is not conserved), if the strong d-pn vertices contain momentum dependent form factors

$$q_\mu T^\mu_{pole} = -\xi_\nu(d)\bar{u}(p')\gamma^\nu[G(-k_t^2) - G(-k_s^2)]C\bar{u}^T(n') \neq 0.$$

Therefore the Born current in eq.(3) is not complete and a *contact* contribution (which should not contain any pole-type singularities!) must be added to provide current conservation on the tree-level. To this end we use *minimal insertion* of the gauge field directly into the d-pn vertex [11], which gives rise to a *contact amplitude*, T^μ_c:

$$T^\mu_c = \xi_\nu(d)\bar{u}(p')\gamma^\nu \int_0^1 \frac{d\lambda}{\lambda}\frac{\partial}{\partial q_\mu}\{e_p G(-(k_s - \lambda q/2)^2) + e_n G(-(k_s + \lambda q/2)^2)\}C\bar{u}^T(n'). \quad (4)$$

Calculating the contraction $q_\mu T^\mu_c$ with $e_p = 1$ and $e_n = 0$, one can easily check that the total current is conserved [10]: $q_\mu(T^\mu_{pole} + T^\mu_c) = 0$, irrespective of the explicit form of the strong form factor $G(-k^2)$, and hence the total amplitude is gauge invariant.

To proceed we need make a choice for the vertex function $G(-k_i^2)$. Previous work [1] showed that in the energy region above $E_\gamma > 1$GeV a conventional potential model wave function cannot describe the data.

In the present approach it is hypothesized that the d-np vertex can be separated into two parts: a *soft* part corresponding to conventional meson exchange theory, and a small *hard* component caused by short-range phenomena. It is assumed that the soft part describes all low-energy (static) properties of the NN system and provides the dominant contribution to the normalization of the bound state wave function, while the hard part dominates the cross section at large virtualities. Since the microscopic structure of the short-range dynamics is poorly known we will use an effective counting rule prescription [13] to describe the hard part of the d-pn vertex:

$$G(p^2) = \frac{C}{(\Lambda^2/2 + m^2 - p^2)^g}, \quad (5)$$

where p is the momentum of the off-shell nucleon, C is a normalization parameter and Λ is related to the inverse of the range. For the special case $g = 3$ [13] the three-pole vertex represents one meson propagator and two (monopole) nucleon-meson form factors.

At the large virtualities involved obviously relativistic effects play an important role. In practice there exist various relativistic formulations, such as the instant form (if) formalism and the light-front (lf) approach. Whereas in an exact calculation these are expected to yield the same result, this is not true in a truncated Fock space scheme. Indeed it has been noted [14,16] that in lowest order (IA) the lf and if approaches lead to different results. In particular in the lowest order Fock states in the lf approach negative energy states do not enter. To illustrate this model dependence in the following we distinguish the covariant (instant form) and the light-front approach.

Covariant Approach: Instant form kinematics

In terms of the variables k_i the vertex in eq.(5) takes on the form ($i = s, t, u$)

$$G^{\rm if}(-k_i^2) = \frac{C}{(\delta^2 - k_i^2)^g}, \tag{6}$$

where $k_t^2 = -\alpha\vec{k}^2$, $k_u^2 = -(2-\alpha)\vec{k}^2$, $k_s^2 = -\vec{k}^2$, with $\vec{k}^2 = s/4 - m^2$ and $\alpha = 2pq/(dq) = 1 - \frac{k}{E_k}\cos\theta$. Furthermore $\delta^2 = \Lambda^2/4 + \alpha_0^2$ with $\alpha_0^2 = m^2 - m_d^2/4$.

Substituting eq.(6) in eq.(3) the cross section can be obtained straightforwardly. The absolute value of the cross section cannot be predicted and therefore the (only) parameter C is fixed by fitting the data at E_γ=1 GeV. We have chosen this energy for the normalization since it is in this region that the microscopic d-np vertex gives a reasonable description of the absolute value of the empirical cross section.

The resulting cross section $s^{11}\frac{d\sigma}{dt}$ at $\theta = 90°$ is compared with data in fig. 1 for $g = 3$ and $g = 4$. For comparison the results for the reduced amplitude approach of ref [7] which is very close to that $g = 4$, are also shown. It is seen that the observed overall energy dependence of the cross section in the energy region $1 < E_\gamma < 4$GeV is described well with $g = 3$. Only at the highest energies a larger value of g would fit better.

Turning to the angular distribution one sees from fig. 2 (dot-dashed curve) that in the if

approach the predicted angular dependence on θ increases rapidly away from 90°. Although the data are available only for a few angles this is clearly not observed experimentally.

Light-front approach

Here the light-cone variables α and $k_\perp$ have a direct physical interpretation as the longitudinal and transverse momentum fraction carried by the nucleon in the deuteron, with $\vec{k}^2 = (s-4m^2)/4 = \frac{m^2(1-\alpha)^2+k_\perp^2}{\alpha(2-\alpha)}$. We will oriente the normal vector of the light-cone hypersurface along $\vec{q}$ to suppress Z-graphs [10]. In general strong lf form factors are functions of two variables, for which we can use any convenient pair from the set $(\alpha, k_\perp, k_3, \vec{k}^2)$. For simplicity we will assume a factorized form of the lf form factor:

$$G^{\mathrm{lc}}(\alpha, \vec{k}^2) = \frac{C}{(\beta^2 - \vec{k}^2)^g}\phi_\kappa(\alpha), \text{ with } \phi_\kappa(\alpha) = \alpha^{-\kappa}. \tag{7}$$

Here the functional dependence on $\vec{k}^2$ is the same as in the instant form, eq.(6). The function $\phi(\alpha)$ in eq.(7) goes to unity in the nonrelativistic limit. The simplest choice is $\kappa = 0$, (as in ref [16]), but a more realistic choice is $\kappa = \frac{1}{2}$ [17] (basically corresponding to the Wick Cutkovsky model). At $\theta = 90^\circ$ (where $\alpha = 1$) both vertices are identical and the lf and if formalisms lead to the same cross cross section (apart from a slight difference coming from the contact term, eq.(4)). However, at $\theta \neq 90^\circ$ the form factors (6) and (7) reproduce a completely different dynamics. Note, that the angular dependence of the cross section arises mainly from the dependence of the arguments of G in eq.(3) on $\alpha = (1 - \frac{k}{E_k}\cos\theta)$.

In G^{if} in eq.(6) these arguments are $k_t^2 = -\alpha\vec{k}^2$ and $k_u^2 = -(2-\alpha)\vec{k}^2$. This gives rise to a strong increase of the cross section at forward and backward angles, which has a dynamic origine. Namely because of the requirement of covariance in the instant form the vertex dependence on α and $\vec{k}^2$ effectively reduces to the dependence on one covariant variable only, say k_t^2 (or k_u^2); the latter is a function of both s and θ, and hence in this case the angle dependence is essentially dictated by the value of the vertex parameter g in eq.(5), which also determines the s- dependence.

On the other hand in the lf approach the angular distribution is flatter since in this case the vertex depends on two variables, α and $\vec{k}^2$, which are in principle independent. Thus,

a steep s-dependence of the cross section may be consistent with a flat θ-dependence. One sees from the curves in fig.2 that with $\kappa = \frac{1}{2}$ the observed angular dependence at $E = 3.2$ GeV is described well.

In fig.3 the calculated energy dependence of the cross sections at 4 different angles is compared with the data. While the overall agreement is reasonable the data at the forward angles suggest a steeper increase of the cross section with energy than predicted, and moreover there is a discrepancy for the highest energy at 69°. Also shown are the results of the RNA [7] and (except at 69°) the QGS approach [8]; the latter is expected to be applicable only at small angles.

III. DISCUSSION

It is of interest to discuss the underlying mechanism for scaling in the present approach. Using eqs.(2-4) we can express the cross section for large (but finite) s

$$\frac{d\sigma}{dt} = \frac{m_d C}{\sqrt{s}(s-m_d^2)^{2g-1}(s-m_d^2)^{3/2}} f(\theta, s). \tag{8}$$

For $\theta = 90°$ one has $\alpha = 1$ and $f(\theta, s) = 1$. In the relevant region of s the contribution of m_d^2 in eq.(8) is still non-negligible; for $g = 3$ the calculated "prescaling" behavior in the relevant energy region can be approximately written as $\frac{d\sigma}{dt} \approx s^{-n+2}$, with $n - 2 = 10$.

For other angles $f(\theta, s)$ depends upon s through $\alpha(\theta, s)$ and therefore is a different function of s at different angles, effectively giving an additional power of s in eq.(8) when $\theta \neq 90°$. Therefore only for $s \to \infty$ the longitudinal fractions α and $\alpha' = 2 - \alpha$ do not depend upon s: $\alpha(s \to \infty) \to 1 - \cos(\theta)$, and one expects an (angular independent) scaling. We find that in the present model in the energy region 1-4 GeV n decreases slowly from $n - 2 \approx 10$ at 90° to $n - 2 \approx 8$ at 10° (or 170°). In practice from eqs.(2-4) this behavior can be expressed as an angular dependence of the cross section on the lf variable α, namely $\frac{d\sigma}{dt} \propto \alpha^{-2}$ ($\alpha > 1$) or $\propto$ $(1-\alpha)^{-2}$ ($\alpha < 1$).

In the past Brodsky [7] has discussed the concept of a "hadron helicity conservation law".

In case of the $\gamma + d \to p + n$ reaction it states that only helicity amplitudes satisfying the condition $\lambda_p + \lambda_n = \lambda_d$ contribute at $s \gg m^2$. Taking into account that only the Dirac coupling ($\approx \gamma_\mu$) in the γNN-vertex conserves hadron helicity, while Pauli coupling ($\approx \sigma_{\mu\nu}$) does not, we see that in the present approach asymptotic "hadron helicity conservation" will occur only in case the Pauli couplings are negligible in the limit $s \gg m^2$, i.e. if the latter fall off at least as $1/s^2$. Note that the gauge constraint for the 3-point EM Green function, which does not allow any form factors in the Dirac coupling (in case of a reducible vertex with real photon), does not lead to any restrictions for the Pauli one.

In practice helicity conservation predicts that for $s \gg m^2$ the cross section in the backward hemisphere receives no contribution from the neutron pole (located near 180°) but only from the deuteron pole which does not depend on the angle. Hence at backward angles the cross section would scale with s independent of angle. On the other hand at forward angles, where we have a competition of two different pole contributions, namely the proton one (which depends on θ and s) and the deuteron one, one does not expect a unified s-dependence but rather an angle dependent scaling. For this reason it would be of interest to extend the experiment to backward angles.

As to the differences in the results in the lf and if formalisms we note that we used an effective counting rule for the hard part of the d-pn vertex; this is based on the assumption of a *fixed number of constituents* which is only well defined in the lf approach, but not in the if. Therefore we consider the results of the latter (in this particular model) less reliable at extreme angles, which involve large t-, u-virtualities.

We note that the deuteron vertex at large virtualities can in principle be addressed in more detail in semi-exclusive reactions, such as $d(e, e'p)X$ at large Q^2 in which the spectator proton is observed in the backward hemisphere to avoid contamination from hadronization products. In the past this reaction has been proposed [15] to discriminate between various models for the nuclear EMC effect. In the IA the deuteron hadronic tensor factorizes in terms of a neutron structure function and a nuclear spectral function determining the variables $\alpha, k_\perp$ of the observed recoiling proton. Hence the cross section is directly proportional to the

square of the deuteron vertex. In this respect we note that the relation between the spectral function and the deuteron wave function is different for the lf and if formalism [15,16]. (In fact in the covariant formulation it is not well possible to define a proper normalization.) This difference tends to increase with increasing α, and depend also on $k_\perp^2$ [18].

Finally we note that in this work we have assumed (as an extreme model) that the hard part of the NN interaction resides only in the (initial) deuteron vertex, and not in the final np state; in future work we will explore the contribution from hard contributions in the final np rescattering.

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Figure Captions

Fig. 1. Calculated energy dependence of $\frac{d\sigma}{dt}s^{11}$ at $\theta = 90°$ for $g = 3$ (dashed), 4 (dotted curve); the solid curve is the result from Brodsky and Hiller [7]; the data are from [4,5].

Fig. 2. Calculated angular dependence of $s^{11}\frac{d\sigma}{dt}$ at $E_\gamma = 3.2$ GeV , for $g = 3$, in the light-front approach with κ=0 (solid),$\frac{1}{2}$ (dotted), 1 (dashed), and instant form (dashed-dotted curve).

Fig. 3. The scaled cross section $s^{11}\frac{d\sigma}{dt}$ as a function of photon energy for (a) $\theta = 89°$, (b) 69°, (c) 52°, (d) 36°. The data are from ref. [4,5], the solid curve shows the present results, the dashed curve those from [8] (not available for 69°), and the dotted curve those from [7].

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